Revisiting cosmological bounds on radiative neutrino lifetime

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Neutrino oscillation experiments and direct bounds on absolute masses constrain neutrino mass differences to fall into the microwave energy range, for most of the allowed parameter space. As a consequence of these recent phenomenological advances, older constraints on radiative neutrino decays based on diffuse background radiations and assuming strongly hierarchical masses in the eV range are now outdated. We thus derive new bounds on the radiative neutrino lifetime using the high precision cosmic microwave background spectral data collected by the Far Infrared Absolute Spectrophotometer instrument on board of Cosmic Background Explorer. The lower bound on the lifetime is between a few$\times 10^{19}$ s and $\sim 5 \times 10^{20}$ s, depending on the neutrino mass ordering and on the absolute mass scale. However, due to phase space limitations, the upper bound in terms of the effective magnetic moment mediating the decay is not better than $\sim 10^{-8}$ Bohr magnetons. We also comment about possible improvements of these limits, by means of recent diffuse infrared photon background data. We compare these bounds with pre-existing limits coming from laboratory or astrophysical arguments. We emphasize the complementarity of our results with others available in the literature.

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I. INTRODUCTION

The last decade has seen two great improvements in astroparticle physics: the wealth of information on neutrino physics (for a recent review see e.g. [1]) and the impressive precision of new cosmological data (see for example the latest WMAP team results [2]). The standard cosmological scenario predicts the existence of a diffuse background of low-energy neutrinos, which has been often investigated in the past to probe non-standard neutrino properties (for a few recent examples, see [3, 4, 5, 6]). At this stage, it is meaningful to reassert the impact of non-standard physics in the neutrino sector on cosmological observables, or equivalently to re-explore the constraints provided by cosmology on exotic physics.

In this paper we revisit the bounds on neutrino radiative lifetime coming from cosmology. Indeed, older constraints based on the diffuse Cosmic Infrared Background (CIB) and assuming strongly hierarchical masses in the eV range [7, 8] (see also [9, 10, 11] for general reviews) are now outdated and strictly speaking inapplicable. The neutrino mass splittings squared provided by oscillation experiments and present upper bounds on the neutrino mass scale constrain neutrino mass differences to fall in the microwave energy range ($E \sim 10^{-3}$ eV), in most of the allowed parameter space. Proper bounds must be derived using the high precision cosmic microwave background (CMB) data collected by the Far Infrared Absolute Spectrophotometer (FIRAS) instrument on board the Cosmic Background Explorer (COBE), which tested the blackbody nature of the spectrum at better than 1 part in 10^4 [12, 13]. The high precision of this measurement has been already exploited to constrain some new physics scenarios, producing deformations on the CMB spectrum [14, 15]. In addition, a new estimate of the CIB flux (although less accurate than CMB) has been recently derived [16] from the SPITZER telescope data [17]. Consequently, one can also infer updated bounds on the radiative lifetime when the unredshifted photon energy falls in this infrared range ($E \sim 10^{-2}$ eV), i.e., in the limit of non-degenerate neutrino mass pattern.

Here we perform such an analysis for radiative neutrino decays, and compare the new bounds obtained with limits coming from laboratory or astrophysical arguments. Although taken at face value these bounds are not competitive with the most stringent ones already available, we emphasize that they probe different combinations and/or regimes of the effective couplings describing the electromagnetic properties of neutrinos, thus being complementary (rather then redundant) with respect to the others available in the literature. The plan of our work is as follows. In Sec. II we summarize the relevant formalism, while in Sec. III we present the CMB data used and the bounds obtained. In Sec. IV we briefly discuss the bounds coming from

the CIB flux. In Sec. V we comment our results and give the conclusion.

II. RADIATIVE NEUTRINO DECAYS

Let us denote by ν_i the (active) neutrino fields respectively of masses m_i, where $i = 1, 2, 3$. The radiative decay $\nu_i \to \nu_j + \gamma$ can be thought of as arising from an *effective* interaction Lagrangian of the form

$$\mathcal{L}_{\rm int} = \frac{1}{2}\bar{\nu}^i \sigma_{\alpha\beta}(\mu_{ij} + \epsilon_{ij}\gamma_5)\nu^j F^{\alpha\beta} + \text{ h.c.} \qquad (1)$$

where $F^{\alpha\beta}$ is the electromagnetic field tensor, $\sigma_{\alpha\beta} = [\gamma_\alpha, \gamma_\beta]$ where γ_μ are the Dirac-matrices and $[.,.]$ is the commutator, ν_i is the neutrino field of mass m_i, and μ_{ij} and ϵ_{ij} are the magnetic and electric transition moments usually expressed in units of the Bohr magneton μ_B. The convention to sum over repeated indices is used. In general, μ_{ij} and ϵ_{ij} are functions of the transferred momentum squared q^2, so that constraints obtained at a different q^2 are independent. The radiative decay rate for a transition $i \to j$ is written

$$\begin{aligned}\Gamma_{ij}^\gamma &= \frac{|\mu_{ij}|^2 + |\epsilon_{ij}|^2}{8\pi}\left(\frac{m_i^2 - m_j^2}{m_i}\right)^3 \\ &\equiv \frac{\kappa_{ij}^2}{8\pi}\left(\frac{m_i^2 - m_j^2}{m_i}\right)^3 . \end{aligned} \qquad (2)$$

In the following, we shall quote the bounds in terms of κ_{ij}^2. We shall assume that the radiative decay rate is very low compared with the expansion rate of the universe; neither the cosmological evolution or the primordial neutrino spectrum is affected by the additional coupling we are going to introduce. A posteriori, this is known to be an excellent approximation. For the same reason, we shall also neglect "multiple decays" (the daughter neutrino ν_j constitutes a negligible fraction of the original ν_i quasi-thermal population). We shall take our input data for neutrino mass eigenstate densities from the calculation performed in [18] without any extra parameter, as non-vanishing chemical potentials. With present data, the latter are anyway constrained to be well below $\mathcal{O}(1)$ [19], so dropping this assumption would not change much our conclusions.

From simple kinematical considerations it follows that in a decay $\nu_i \to \nu_j + \gamma$ from a state of mass m_i into one of mass $m_j < m_i$, the photon in the rest frame of the decaying neutrinos is thus monochromatic (two-body decay), with an energy

$$\varepsilon_{ij} = \frac{m_i^2 - m_j^2}{2\, m_i} . \qquad (3)$$

At present, the neutrino mass spectrum is constrained by the well-known values of the two squared mass splittings for the atmospheric (Δm_H^2) and the solar (Δm_L^2) neutrino problems. We take their best-fit values and 2σ ranges from [1]:

$$\begin{aligned}\Delta m_L^2 &= 7.92\,(1 \pm 0.09) \times 10^{-5}\ \text{eV}^2 , \qquad (4)\\ \Delta m_H^2 &= 2.6\,(1^{+0.14}_{-0.15}) \times 10^{-3}\ \text{eV}^2 . \qquad (5)\end{aligned}$$

The remaining unknowns in the neutrino spectrum are the absolute mass scale (equivalently, the mass of the lightest eigenstate m_1) and the mass hierarchy. Namely, in normal hierarchy (NH) the mass pattern would be

$$\begin{aligned} & m_1 \ , \\ m_2 &= \sqrt{m_1^2 + \Delta m_L^2}\ , \\ m_3 &= \sqrt{m_1^2 + \Delta m_L^2 + \Delta m_H^2}\ ; \end{aligned} \qquad (6)$$

while in inverted hierarchy (IH) one would have

$$\begin{aligned} & m_1 \ , \\ m_2 &= \sqrt{m_1^2 + \Delta m_H^2}\ , \\ m_3 &= \sqrt{m_1^2 + \Delta m_L^2 + \Delta m_H^2}\ . \end{aligned} \qquad (7)$$

In the limiting case of normal hierarchy and $m_1 = 0$, the lightest neutrino for which a decay is possible has a mass $m_2 \simeq 9 \times 10^{-3}$ eV and is thus non-relativistic for most of the universe lifetime, namely in the redshift range $z \lesssim 50$. We can thus safely work in the approximation of all neutrinos decaying effectively at rest. In this limit, we can also neglect the momentum distribution of the neutrino spectra. The formalism which would allow one to generalize our results to the momentum-dependent case has been developed in [20], which we address the interested reader for further details. However, the corrections are small, of the order of powers of the neutrino temperature to mass ratios, and also vanishing in the limit of very long lifetimes.

We shall discuss the limits on κ_{ij}^2 as a function of m_1 and for the two cases NH and IH. We shall vary the mass scale in $0 \lesssim m_1 \lesssim 2$ eV as allowed by the Mainz experiment on the ^{3}H beta decay endpoint [21]. In this respect, we shall be conservative: If neutrinos are Majorana particles the more stringent bound from $0\nu\beta\beta$ searches apply, with an effective mass bound $m_{\beta\beta} < 0.81$ eV [22]. Structure formation, combined with other cosmological data, also constrains $\sum_i m_i$ [2]. Present cosmological bounds span the range $\sum_i m_i \lesssim 0.2 - 2$ eV, [2, 22, 23, 24, 25, 26] depending on the data sets used and priors assumed. An upper limit of $\sum_i m_i \sim 0.6$ eV (i.e. $m_1 \simeq 0.2$ eV) is often considered robust, and we shall report it for illustrative purposes. Yet, one may circumvent the

$0\nu\beta\beta$ bound (e.g. with a Dirac neutrino) and significantly relax the most stringent cosmological bounds (for example with a conservative combination of cosmological data sets and priors or with an exotic dark energy sector), so in the following we shall present our results up to the value $m_1 = 2\,\mathrm{eV}$.

Let F_E be the present energy flux of photons with present energy E produced by neutrino decay. The differential energy flux φ_E (energy flux F_E per unit energy and solid angle) is related to the differential number flux φ_n (the particle flux F_n per unit energy and solid angle) at present by

$$\varphi_E \equiv \frac{\mathrm{d}^2 F_E}{\mathrm{d}E\,\mathrm{d}\Omega} = E\,\frac{\mathrm{d}^2 F_n}{\mathrm{d}E\,\mathrm{d}\Omega} \equiv E\,\varphi_n, \tag{8}$$

and it can be shown that, if the lifetime τ_i of the neutrino of mass m_i is much greater than the universe lifetime it holds [20] [1]

$$\varphi_E = \frac{\Gamma_{32}^{\gamma}}{4\pi}\frac{n_3}{H(z_{32})} + \frac{\Gamma_{31}^{\gamma}}{4\pi}\frac{n_3}{H(z_{31})} + \frac{\Gamma_{21}^{\gamma}}{4\pi}\frac{n_2}{H(z_{21})}\,, \tag{9}$$

where $n_i \simeq 113\ \mathrm{cm}^{-3}$ is the present number density of the $i-$th neutrino *in absence of decay*, the Hubble function is (assuming, for simplicity, a flat cosmology) $H(z) = H_0\sqrt{\Omega_M(1+z)^3 + \Omega_\Lambda}$, $H_0 \simeq 73$ km s^{-1} Mpc^{-1} being the present Hubble expansion rate, and $\Omega_M \simeq 0.26$ and $\Omega_\Lambda \simeq 0.74$ respectively the matter and the cosmological constant energy density relative to the critical one. The dependence on energy enters implicitly via the quantities $0 \leq z_{ij} = \varepsilon_{ij}/E - 1$.

In practice, to a very good approximation one can write a general equation of the kind

$$\varphi_E = \frac{\Gamma_{H}^{\gamma}}{4\pi}\frac{n_H}{H(z_H)} + \frac{\Gamma_{L}^{\gamma}}{4\pi}\frac{n_L}{H(z_L)}\,, \tag{10}$$

where, however, the meaning of the factors however depends on the hierarchy. In NH, in the first two terms of the sum in Eq. (9) it holds $z_{32} \simeq z_{31} \equiv z_H$, and one can identify $z_L = z_{21}$, $\Gamma_L^\gamma = \Gamma_{21}^\gamma$, $\Gamma_H^\gamma = \Gamma_{31}^\gamma + \Gamma_{32}^\gamma$. In IH, it is the last two terms of the sum in Eq. (9) which have $z_{31} \simeq z_{21} \equiv z_H$, and using $n_2 \simeq n_3$ one can identify $z_L = z_{32}$, $\Gamma_L^\gamma = \Gamma_{32}^\gamma$, $\Gamma_H^\gamma \equiv \Gamma_{31}^\gamma + \Gamma_{21}^\gamma$. In both cases, we shall therefore express our bounds in terms of $\kappa^2_{L,H}$ keeping in mind their slightly different meaning for the two cases of NH and IH.

In Fig. 1 we represent the unredshifted photon energy ε_{ij} from decaying neutrinos [Eq. (3)] as a function of the lightest neutrino mass eigenstate

[1] Note that τ_i may be much shorter than the radiative lifetime, which in most exotic models is dominated by invisible decays [4, 27, 28]. In the present work we are neglecting these cases.

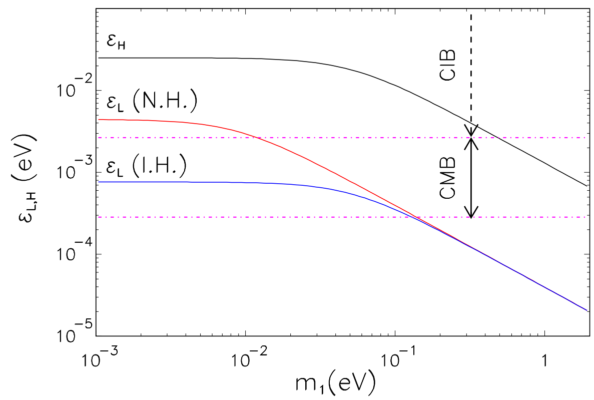

FIG. 1: Unredshifted photon energy ε from decaying neutrinos [Eq. (3)] as a function of the lightest neutrino mass eigenstate m_1, for the two neutrino mass splittings (L,H) in normal and inverted hierarchy. (See text for details) The horizontal band represents the energy range of the CMB spectrum measured by FIRAS [12]. The CIB energy range is also shown.

m_1 in the case of normal and inverted hierarchy, where the meaning of $\varepsilon_{L,H}$ is clear from the previous discussion. We also indicate by an horizontal band the energy range of the CMB spectrum ($2.84 \times 10^{-4}\,\mathrm{eV} \leq E \leq 2.65 \times 10^{-3}\,\mathrm{eV}$) measured by FIRAS [12]. We also show the CIB range in the energy band above the FIRAS range up to (conventionally) 0.15 eV [16]. For $m_1 \lesssim 0.5$ eV, the photon energy ε_H falls in the CIB range.

For photons emitted at $z = 0$ in the FIRAS range, the effect of radiative decays is most prominent and results in a feature on the CMB spectrum. Actually even if photons are emitted at higher energy the effect is still strong, since photons emitted at a redshift of a few enter the FIRAS spectrum because of cosmological redshift; it is easy to check that one has thus some sensitivity to κ_H in the whole range for m_1. However, as we will see, for $m_1 \lesssim 0.1$ eV a stronger (but less robust) limit can be obtained using directly the CIB data.

On the other hand, for $m_1 \gtrsim 0.14$ eV the photons corresponding to the smaller splitting are falling in the radio band, below the frequency range probed by COBE, where measurements are more uncertain and thus one has no sensitivity to κ_L and the corresponding bound disappears.

III. THE CMB BOUND

To constrain the neutrino electromagnetic decay we use the COBE/FIRAS data for the experimentally measured CMB spectrum, corrected for foregrounds [12]. Note that the new calibration of FIRAS [13] is within the old errors and would not change any of our conclusions. The $N = 43$ data points Φ_i^{exp} at different energies E_i are obtained by

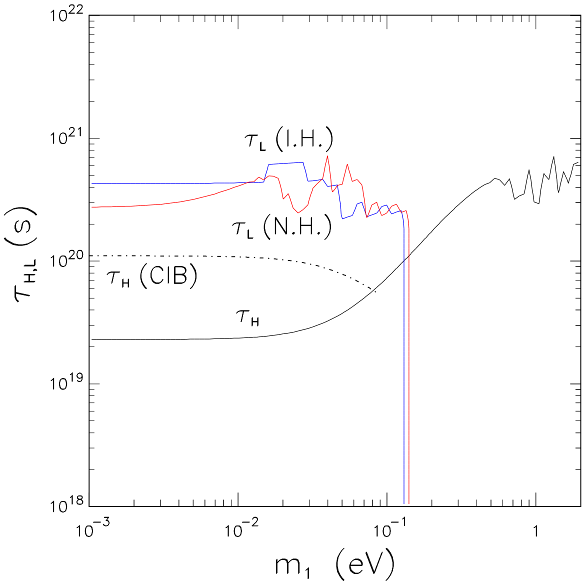

FIG. 2: Bounds on τ_H and τ_L vs. m_1, for the two cases of NH and IH. The regions below the solid curves are excluded at 95 % C.L. The curves for τ_H coincide in the two cases, although the definition of τ_H is different (see text). The dot-dashed line represents the limit on τ_H obtained from cosmic infrared background.

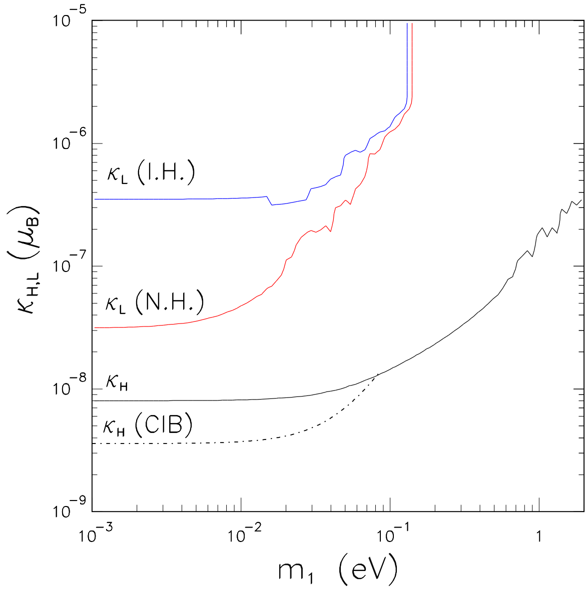

FIG. 3: Bounds on κ_H and κ_L vs. m_1, for the two cases of NH and IH. The regions above the solid curves are excluded at 95 % C.L. The curves for κ_H coincide in the two cases, although the definition of κ_H is different (see text). The dot-dashed line represents the limit on κ_H obtained from cosmic infrared background.

summing the best-fit blackbody spectrum to the residuals reported in Ref. [12]. The experimental errors σ_i and the correlation indices ρ_{ij} between different energies are also available. In the presence of neutrino decay, the original radiance (energy flux per unit of solid angle) of the "theoretical blackbody" at temperature T

$$\Phi^0(E,T) = \frac{E^3}{4\pi^3}\left[\exp(E/T)-1\right]^{-1} \qquad (11)$$

would gain an additional term so that the intensity becomes

$$\begin{aligned}\Phi^0(E,T) &\rightarrow \Phi(E,T,\kappa^2_{L,H},m_1) \\ &= \Phi^0(E,T)+\varphi_E(\kappa^2_{L,H},m_1). \quad (12)\end{aligned}$$

We then build the reduced chi-squared function

$$\chi^2_\nu(T,\kappa^2_{L,H},m_1) = \frac{1}{N-1}\sum_{i,j=1}^{N}\Delta\Phi_i(\sigma^2)^{-1}_{ij}\Delta\Phi_j\,, \qquad (13)$$

where

$$\Delta\Phi_i = \Phi_i^{\rm exp} - \Phi(E_i,T) \qquad (14)$$

is the i-th residual, and

$$\sigma^2_{ij} = \rho_{ij}\sigma_i\sigma_j \qquad (15)$$

is the covariance matrix. In principle, the parameter T entering initially in $\Phi^0(E,T)$ needs not to be fixed at the standard value $T_0 = 2.725 \pm 0.002$ K [13], which is the best fit of the "distorted" spectrum eventually observed now. The initial T before a significant fraction of neutrinos decays should thus be a free parameter, to be determined in the minimization procedure. Practically, however, the distortion introduced by the neutrino decay spectrum is such highly non-thermal that a change in T can not be accomodated for any significant neutrino lifetime: the constraints obtained fixing T are basically the same.

Our results are reported in Fig. 2, where we represent the exclusion plot in the plane $\tau_{L,H} \equiv (\Gamma^\gamma_{H,L})^{-1}$ vs. m_1, where the regions below the solid curves are excluded at 95 % C.L. For small values of m_1 the most stringent limit is $\tau_L \gtrsim 4\times10^{20}$ s in IH (slightly better than in NH case), while the bound on τ_H is about an order of magnitude smaller, say $\tau_H \gtrsim 2\times10^{19}$ s, since for low m_1 only photons produced by H decays at a redshift of few are in FIRAS range. On the contrary, for $m_1 \gtrsim 0.14$ eV, the bound on τ_L disappears, while the bound on τ_H becomes more stringent, being $\tau_H \gtrsim 5\times10^{20}$ s. Note that the "fuzzy" behaviour of the bounds is due to the sharp edge of the photon spectrum at $E = \varepsilon_{H,L}$: when the photon energy embeds a new FIRAS bin, the χ^2 function has a sharp discontinuity.

In Fig. 3 we translate the plot of Fig. 2 in an exclusion plot in the plane m_1 vs. $\kappa_{L,H}$. Here the factor $(\delta m^2_{ij}/m_i)^3$ maps in a non trivial way the bounds in terms of $\kappa_{L,H}$. The regions above the solid curves are excluded at 95 % C.L. For the NH case, $\kappa_L \lesssim 3\times10^{-8}\,\mu_B$, while in the IH

case, $\kappa_L \lesssim 3 \times 10^{-7}\,\mu_B$. In agreement with our previous considerations, the bound on κ_L disappears for $m_1 \gtrsim 0.14\,\mathrm{eV}$. On the contrary, the bound for κ_H is always present, and it corresponds to $\kappa_H \lesssim 8 \times 10^{-9}\,\mu_B$ apart for the degenerate region, where it degrades down to $10^{-7}\,\mu_B$ or even more. Note also that typical cosmological upper bounds would already exclude the extreme degenerate case.

IV. THE CIB BOUND

The Cosmic Infrared Background (CIB) is mainly the relic emission of the formation and evolution of the galaxies of all types at wavelengths larger than a few microns. The spectrum of the CIB is peaked around $\sim$ 100 μm ($E \sim 1.2 \times 10^{-2}$ eV), thus just is in the energy range ε_H of photon from radiative ν decays, for $m_1 < 0.1$ eV. Recently, a new estimate of the CIB flux has been established using the Spitzer Observatory data [16]. The measured CIB flux is $\Phi_{\rm CIB} \sim 24$ nW m^{-2} sr^{-1}. Using this number we can obtain a rough bound on τ_H (and hence on κ_H) simply requiring that the total energy flux of the photons coming from ν decay does not exceed the CIB flux:

$$\int_{E_{\rm min}}^{\varepsilon_H} \varphi_E \, dE < \Phi_{\rm CIB} \ , \tag{16}$$

where we consider as lower limit of the CIB range the upper value of the FIRAS range, i.e. $E_{\rm min} = 2.65 \times 10^{-3}$ eV. The bounds of τ_H and κ_H obtained from Eq. (16) are shown respectively in Figs. 2 and 3 by the dot-dashed line [2]. Although these bounds are stronger than those obtained by the FIRAS data in the same range of m_1, we emphasize that they should be considered only as indicative. In fact, the CIB flux have still strong uncertainties and the precise spectral shape is essentially unknown (a factor ~ 3 of uncertainty should be accounted [16]).

It is interesting to comment that, if it turns out that $m_1 \lesssim 0.1\,\mathrm{eV}$, an improvement on the bound on $\tau_{\rm H}$ will clearly take advantage of a better measurement of the CIB flux and a more detailed knowledge of the astrophysical sources contributing to it. Conversely, $m_1 \gtrsim 0.1\,\mathrm{eV}$ would imply a significant degree of neutrino clustering in large dark matter halos, the larger the mass the stronger the clustering [29, 30]. In turn, the expectation of overdensities would motivate analyses in the microwave sky toward specific targets (like nearby galaxies or galaxy clusters), thus taking full advantage of the spectral feature expected from neutrino decay and looking for an angular-dependent enhancement over the CMB background. In spirit, this would be similar to what performed in the X-ray band when searching for signatures of sterile neutrino or axion decays (see e.g. [31, 32]). Although a detailed treatment of these issues goes beyond the purpose of this work, it is worth noting that which one is the regime to consider will be basically answered by the KATRIN experiment on tritium beta decay [33].

V. DISCUSSION AND CONCLUSIONS

In this paper, we have revisited the bounds on the neutrino radiative lifetime coming from cosmology, deriving updated constraints from the high precision CMB spectrum data collected by the FIRAS instrument on board of COBE. We also compare these bounds with those obtained (by a back-of-the-envelope calculation) using the measurement of the flux of the Cosmic Infrared Background, which although sometimes overrides those coming from the CMB data, should be considered only as qualitative.

Previous cosmological bounds were either derived in a pre-COBE era [7] or there were assumed higher masses and/or a strongly hierarchical mass spectrum, thus using the infrared background to derive the constraints [8]. This has motivate us to re-evaluate the bounds within the presently allowed range of parameters suggested by neutrino oscillation physics and tritium endpoint experiments. Since it is customary to parameterize the neutrino electromagnetic decay via an effective operator of the kind reported in Eq. (1), it makes sense to translate the bounds (which actually are on the lifetime) into bounds on the parameters $\kappa_{H,L}$ [see Eqs. (2,10)]. These constraints are not better than $\kappa \lesssim 10^{-8}\,\mu_B$ which at first sight do not appear competitive with astrophysical limits neither with most of the laboratory bounds [34]. Nevertheless, every experimental measure and every cosmological and astrophysical constraint has its own systematic uncertainties and its own recognized or un-recognized loop-holes. In this sense it is certainly important to use many different approaches to constrain fundamental neutrino properties.

In particular, the cosmological bound is based on the appearance of the daughter photons, and thus is very direct (modulo the underlying cosmological assumptions). Therefore the bounds on $\tau_{H,L}$ are completely independent from the underlying model that mediates the neutrino de-

[2] For NH and $m_1 \lesssim 10^{-2}$ eV also ε_L falls marginally in the CIB range. However, we have explicitly checked that the FIRAS constraint on ε_L is always stronger.

cay. Each decay model must face with our direct bounds on $\tau_{H,L}$ wich are the strongest attainable with direct livetime measurements. Moreover, it also probes the energy scale $q^2 \lesssim 10^{-3}\,\mathrm{eV}^2$, unaccessible to both laboratory experiments and stellar arguments. Other constraints have different features. Neutrino electromagnetic couplings are actually tightly constrained by energy-loss arguments in stars, in particular via the plasmon process $\gamma^* \to \bar{\nu}\nu$, to be $\kappa_{ij} \lesssim 3\times 10^{-12}\,\mu_B$ [35, 36]. These are indirect limits, which strictly speaking apply to $q^2 \gtrsim \mathrm{keV}^2$. Laboratory bounds are obtained via elastic $\nu\, e$ scattering, where the scattered neutrino is not observed, and are at most at the level of $10^{-10}\,\mu_B$ for the electron flavor [34, 37]. The combinations of matrix elements ϵ_{ij} and μ_{ij} that are constrained by various experiments depend on the initial neutrino flavor and on its propagation between source and detector. Cancellations may occur in exotic cases [10], and in any case the energy scale probed is $q^2 \gtrsim \mathrm{MeV}^2$. An additional motivation for independent checks is that models with a strong dependence of $\kappa_{ij}(q^2)$ have been proposed [38]. Our bound seems to exclude extreme runnings of the effective couplings. Vice versa, if one assumes that the elements $\epsilon_{i,j}$ and μ_{ij} are quasi energy independent, laboratory and astrophysical arguments exclude any possibility of phenomenologically interesting radiative decays in cosmology; this is a priori surprising, given that cosmology involves the longest time intervals available and the lowest boosting factors for neutrinos. On the other hand, note that the cosmological bound may be easily violated by invoking a neutrino invisible decay much faster than the universe lifetime. It is intriguing to notice that, despite the impossibility to test this "nightmare case" in laboratory, cosmology may probe to some extent those scenarios, potentially with important consequences for particle physics as well [27, 28].

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